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We extend the previous analysis of the current convective instability as applied to the diffuse auroral situation (Ossakow and Chaturvedi, 1979) to include ion inertial effects, important at high altitudes; and to the case in which ions are highly collisional (non-magnetized), a situation which is realized at E-region altitudes, In the inertial domain the instability growth rates are comparable to those found in the collision dominated domain. This extends the applicability of the instability process to high altitudes. The relevance of the instability to the E-region is dicussed. Finally, electromagnetic effects, which can be important for long wavelengths, are considered and are found to be small in ionospheric situations.

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THE CURRENT CONVECTIVE INSTABILITY AS APPLIED TO THE AURORAL IONOSPHERE

I. INTRODUCTION

Recently the current convective instability has been suggested as a possible mechanism responsible for irregularities causing scintillation phenomena observed by the DNA Wideband Satellite during diffuse auroral situations [Ossakow and Chaturvedi, 1979; Chaturvedi and Ossakow, 1979; Vickrey et al., 1980]. The observations seem to suggest the scintillations as being caused by L-shell aligned, field aligned, large scale irregularities co-located in the regions of soft particle precipitation and horizontal plasma density gradients with the source region confined in latitude [Fremouw et al., 1977; Rino et al., 1978; Rino and Owen, 1980]. The current convective instability occurs in regions where the field-aligned currents exceed a certain threshold, determined by the ambient parameters, etc., in the presence of a plasma density gradient transverse to the magnetic field [Kadomtsev and Nedospasov, 1962].

The linear theory of the instability as applied to the auroral situation appears to give plausible growth times, regions of occurrence, etc. [Ossakow and Chaturvedi, 1979; Vickrey et al., 1980]. Some of the observed features, like L-shell alignment and percentage fluctuation of irregularity amplitudes, may be explained by a nonlinear saturation mechanism of the instability that invokes mode coupling to damped modes in a two-dimensional treatment perpendicular to the magnetic field [Chaturvedi and Ossakow, 1979]. The field-aligned currents cause weak shear in the magnetic field with shear lengths comparable in order of magnitude to the parallel scale sizes of the instability-driven unstable modes. The effect of shear is to reduce the growth rates somewhat but to spatially localize the mode in the density gradient direction [Huba and Ossakow, 1980]. The first global numerical simulations [Keskinen et al., 1980] performed for the nonlinear current

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convective instability show that during the nonlinear evolution the regions of enhancements move equatorward while those of depletions drift polewards. In addition, the power spectrum of the density fluctuations follows a power law with a typical value in the vicinity of $\sim k^{-2}$.

In the present work we have extended the previous linear analysis of the current convective instability [Ossakow and Chaturvedi, 1979] to include ion inertial effects, which assume importance at high altitudes where the ion-neutral collision frequency gets smaller. We find the growth rates to remain of the same order of magnitude in the inertial case as they were in the collisional case [Ossakow and Chaturvedi, 1979]. At the same time, we present the growth rate expression of the current convective instability for the other limit when ions are strongly collisional, i.e., $\nu_{in} \gg \Omega_i$, as is the case in the E-region ionosphere. Further, as is well known, in considering long wavelengths such that $ck/\omega_e \leqslant 1$ (where c is the velocity of light, k the wave number and ω_e the electron plasma frequency), one should consider electromagnetic effects also [Kaw et al., 1974]. We find these effects to be very small in the ionospheric situation.

In the next sections, we first present the growth rate calculations for the ion-inertia dominated regime and for the case of strongly collisional ions, then the effect of electromagnetic corrections on the growth rate, and in the final section we discuss them for the ionospheric application.

II. THEORY

A. Ion Inertia Effects

We shall work in a coordinate system which has its z-axis aligned with the magnetic field, $\underline{B}_o$, with its y-axis pointing northwards, and the x-axis pointing westward. A field-aligned current, along the z-axis and a plasma

density gradient along the y-axis are assumed as a part of the equilibrium. This neglects a weak altitude dependent gradient, any east-west gradient and any equilibrium electric fields. The temperature effects responsible for diffusive damping are ignored for simplicity as is magnetic field shear. The electron inertia is neglected for the low frequencies concerned in the paper and ν_{en}, the electron-neutral collision frequency, is neglected in comparison to ν_{ei}, the electron-ion Coulomb collision frequency, which is a valid assumption at F-region altitudes. The zero-order current is assumed to be caused by an equilibrium drift v_o of electrons over ions along the z-axis. The perturbations are assumed to vary as $\propto \exp\,(ik_x x + ik_y y + ik_z z - i\omega t)$ and the local approximation is used in the analysis which is justified as long as we consider wavelengths to be small compared to the density gradient scale-length ($k_y \gg d \ln n_o/dy$).

With these assumptions the electron velocities are given as

$$v_{ze} \simeq - \frac{eE_z}{m\nu_{ei}} \simeq i \frac{k_z e}{m\nu_{ei}} \tilde{\phi} \tag{1}$$

$$\underline{v}_{\perp e} \simeq \frac{c\underline{E}_{\perp} x \hat{z}}{B_o} \simeq - \frac{c\nabla_{\perp}\tilde{\phi} x \hat{z}}{B_o} \tag{2}$$

where the electrostatic assumptions for perturbed electric fields, $\underline{E} \simeq - \nabla\tilde{\phi}$ has been used.

The ion momentum transfer equation is

$$Mn \frac{\partial \underline{v}_{\perp i}}{\partial t} \simeq en\underline{E}_{\perp} + \frac{en}{c} (\underline{v}_{\perp i} \times \underline{B}_o)_{\perp} - Mn\nu_{in} \underline{v}_{\perp i} \tag{3}$$

The perturbed ion-velocities in the perturbed fields can be written as

$$\underline{v}_{\perp i} \simeq - \frac{1}{\left(1 + \frac{\tilde{\nu}_{in}^2}{\Omega_i^2}\right)} \frac{\tilde{\nu}_{in}}{\Omega_i} \frac{e\nabla_{\perp}\tilde{\phi}}{M\,\Omega_i} - \frac{1}{\left(1 + \frac{\tilde{\nu}_{in}^2}{\Omega_i^2}\right)} \frac{e\nabla_{\perp}\tilde{\phi} x \hat{z}}{M\,\Omega_i} \tag{4}$$

where $\tilde{\nu}_{in} = \nu_{in} - i\omega$, $\Omega_i = \frac{eB_o}{Mc}$. (5)

It is easily verified that under the present approximations

$$\underline{v}_{zi} \simeq 0. \tag{6}$$

The continuity equation for the two species is

$$\frac{\partial n_\alpha}{\partial t} + \nabla_\perp \cdot (n_\alpha \underline{v}_\alpha) + \frac{\partial}{\partial z}(n_\alpha v_{z\alpha}) = 0 \tag{7}$$

with $\alpha \equiv e, i$. Any production and loss due to recombination has been neglected.

Equations (1), (2) and (7) lead to

$$\tilde{\omega}\,\frac{\tilde{n}_e}{n_o} \approx \frac{\partial}{\partial z}\left(\frac{k_z e}{m\nu_{ei}}\,\tilde{\phi}\right) + i\,\frac{e\nabla_\perp \tilde{\phi} x \hat{z} \cdot \underline{\varepsilon}}{m\,\Omega_e}$$

where

$$\tilde{\omega} = \omega - k_z v_o,\ \underline{\varepsilon} = \frac{1}{n_o}\frac{dn_o}{dy}\,\hat{y}. \tag{8}$$

Similarly, eqs. (4), (6) and (7) lead to

$$\omega\,\frac{\tilde{n}_i}{n_o} = \nabla_\perp \cdot \left[i\,\frac{\tilde{\nu}_{in}/\Omega_i}{1 + \frac{\tilde{\nu}_{in}^2}{\Omega_i^2}}\;\frac{e\nabla_\perp\tilde{\phi}}{M\,\Omega_i}\right] +$$

$$\left[i\,\frac{\tilde{\nu}_{in}}{\Omega_i\left(1+\frac{\tilde{\nu}_{in}^2}{\Omega_i^2}\right)}\;\frac{e\nabla_\perp\tilde{\phi}}{M\,\Omega_i} + i\,\frac{e\nabla_\perp\tilde{\phi} x \hat{z}}{M\,\Omega_i\left(1+\frac{\tilde{\nu}_{in}^2}{\Omega_i^2}\right)}\right] \cdot \underline{\varepsilon} \tag{9}$$

The difference of eqs. (8) and (9) leads to

$$v_o k_z \frac{\tilde{n}}{n_o} \approx \nabla_\perp \cdot \left[i \; \frac{\tilde{\nu}_{in}}{\Omega_i} \frac{e\nabla_\perp \tilde{\phi}}{M\,\Omega_i} \right] + i \frac{\tilde{\nu}_{in}{}^2}{\Omega_i{}^2} \; \frac{e\nabla_\perp \tilde{\phi} \; x\hat{z} \cdot \underline{\varepsilon}}{M\,\Omega_i}$$

$$- \frac{\partial}{\partial z} \left(\frac{ek_z}{m\nu_{ei}} \tilde{\phi} \right) \qquad (10)$$

where $\Omega_i >> \tilde{\nu}_{in}$ has been assumed in addition to the quasi-neutrality assumption $\tilde{n}_e \approx \tilde{n}_i = \tilde{n}$.

The dispersion relation is obtained from (8) and (10) and is

$$\omega^2 + i\omega \left(\nu_{in} + \frac{k_z{}^2}{k_\perp{}^2} \frac{\Omega_e \Omega_i}{\nu_{ei}} \right) + \frac{v_o k_z \Omega_i}{k_\perp} \varepsilon \simeq 0 \qquad (11)$$

where some small terms have been neglected. Splitting $\omega = \omega_r + i\gamma$, the growth rate is given as

$$\gamma \simeq - \frac{1}{2} \bar{\nu}_{in} \pm \frac{1}{2} \left(\bar{\nu}_{in}{}^2 + \frac{4 v_o k_z \Omega_i \varepsilon}{k_\perp} \right)^{1/2} \qquad (12)$$

where

$$\bar{\nu}_{in} \equiv \left(\nu_{in} + \frac{k_z{}^2}{k_\perp{}^2} \frac{\Omega_e \Omega_i}{\nu_{ei}} \right)$$

Equation (12) can be maximized with respect to $k_z/k_\perp$ and this maximum growth rate is obtained from the cubic equation for $k_z/k_\perp$

$$\left(\frac{k_z}{k_\perp} \right)^3 - \left(\frac{k_z}{k_\perp} \right) \frac{\nu_{ei} \nu_{in}}{\Omega_e \Omega_i} - \frac{1}{2} v_o \; \frac{\varepsilon \nu_{ei}{}^2}{\Omega_e{}^2 \Omega_i} = 0. \qquad (12a)$$

It is simple to recover the collisional growth rate from (12), for $\bar{\nu}^2{}_{in} >> \frac{4 v_o k_z \Omega_i \varepsilon}{k_\perp}$, one gets

$$\gamma \simeq \frac{\dfrac{V_o k_z}{k_\perp} \dfrac{\nu_{ei}}{\Omega_e} \epsilon}{\dfrac{k_z^2}{k_\perp^2} + \dfrac{\nu_{ei}\nu_{in}}{\Omega_e \Omega_i}} \tag{13}$$

which is the same as before [Ossakow and Chaturvedi, 1979; Chaturvedi and Ossakow, 1979]. The maximum growth rate in this case occurs for (which also results from (12a) by neglecting the last term on the lefthand side of (12a))

$$\left(\frac{k_z}{k_\perp}\right)^2 \simeq \frac{\nu_{ei}\nu_{in}}{\Omega_e \Omega_i}$$

and is

$$\gamma_{CM} \simeq \frac{V_o \epsilon}{2} \left(\frac{m}{M} \frac{\nu_{ei}}{\nu_{in}}\right)^{1/2} \tag{14}$$

In the limit $\nu_{in} \simeq 0$, we recover the purely inertial case from the growth rate expression given by (12),

$$\gamma \simeq -\frac{1}{2} \frac{k_z^2}{k_\perp^2} \frac{\Omega_e \Omega_i}{\nu_{ei}} + \frac{1}{2} \left[\frac{k_z^4}{k_\perp^4} \frac{\Omega_e^2 \Omega_i^2}{\nu_{ei}^2} + \frac{4 V_o \epsilon k_z \Omega_i}{k_\perp}\right]^{1/2} \tag{15}$$

A simplified expression of growth rate for this case can be obtained when

$$1 \gg \frac{4 V_o k_z}{k_\perp} \left(\frac{k_\perp}{k_z}\right)^4 \frac{\nu_{ei}^2}{\Omega_e^2 \Omega_i} \epsilon$$

then one gets

$$\gamma \simeq V_o \epsilon \frac{k_\perp}{k_z} \frac{\nu_{ei}}{\Omega_e}.$$

Maximum growth rate in this case occurs for (from (15) and also results from (12a) by neglecting the second term on the lefthand side of (12a))

$$\frac{k_z}{k_\perp} \simeq \left(\frac{v_o \nu_{ei}{}^2 \varepsilon}{2\Omega_e^2 \Omega_i} \right)^{1/3}$$

and is

$$\gamma_{IM} \approx \left(\frac{v_o \varepsilon}{2} \right)^{2/3} \left(\frac{m}{M} \nu_{ei} \right)^{1/3} \tag{16}$$

It should be noted that in the inertial case the value of $k_z/k_\perp$ which maximizes the growth rate depends on instability driving parameters, v_o and ε, and not just background ionospheric parameters, ν_{ei}, ν_{in}, Ω_e, and Ω_i, as well the case in the collision dominated regime.

B. Collisional *Ions*

We now consider a situation wherein ions are highly collisional such that $\nu_{in} \gg \Omega_i$. This situation is frequently realized at E-region altitudes in ionosphere. The ion and electron velocities may be written as

$$\underline{v}_{\perp i} \simeq \frac{e\underline{E}_\perp}{M\nu_{in}}, \quad v_{zi} \simeq \frac{eE_z}{M\nu_{in}} \tag{17}$$

$$\underline{v}_{\perp e} \simeq \frac{c\underline{E}_\perp x\hat{z}}{B_o}, \quad v_{ze} \simeq - \frac{eE_z}{m\nu_{en}} \tag{18}$$

Upon using the electrostatic assumption, one can readily write the ion and electron continuity equations as

$$\omega \frac{\tilde{n}}{n_o} \simeq - i \frac{ek_\perp{}^2}{M\nu_{in}} \tilde{\phi} \tag{19}$$

and

$$(\omega - v_o k_z) \frac{\tilde{n}}{n_o} \simeq - \frac{c}{B_o} \frac{\underline{k}_\perp x\hat{z} \cdot \nabla n_o}{n_o} \tilde{\phi} + i \frac{ek_z{}^2}{m\nu_{en}} \tilde{\phi} \tag{20}$$

where the quasi-neutrality assumption has been used. The difference of the

above two equations yields

$$v_o k_z \frac{\tilde{n}}{n_o} \simeq -i \frac{ek_\perp^2}{M\nu_{in}} \tilde{\phi} - \frac{ck_x}{B_o L} \tilde{\phi} - i \frac{ek_z^2}{m\nu_{en}} \tilde{\phi} \tag{21}$$

The dispersion relation can be obtained from eqs. (19) and (21),

$$\omega \simeq \frac{v_o k_z}{\left(1 + \frac{M}{m}\frac{\nu_{in}}{\nu_{en}}\frac{k_z^2}{k_\perp^2} - i \frac{\nu_{in}}{\Omega_i}\frac{1}{k_\perp L}\right)} \tag{22}$$

One puts $\omega = \omega_R + i\gamma$, then

$$\omega_R \simeq v_o k_z \left(1 + \frac{M}{m}\frac{\nu_{in}}{\nu_{en}}\frac{k_z^2}{k_\perp^2}\right)^{-1} \tag{22a}$$

The growth rate expression from eq. (22) is

$$\gamma \simeq \frac{v_o k_z \left(\frac{\nu_{in}}{\Omega_i}\frac{1}{k_\perp L}\right)}{\left(1 + \frac{M}{m}\frac{\nu_{in}}{\nu_{en}}\frac{k_z^2}{k_\perp^2}\right)^2 + \left(\frac{\nu_{in}}{\Omega_i}\frac{1}{k_\perp L}\right)^2} \tag{22b}$$

For typical E-region situations, $m/M \sim 2 \times 10^{-5}$, $\nu_{in} \sim 2.5 \times 10^3 s^{-1}$, $\nu_{en} \sim 4 \times 10^4 s^{-1}$, and assuming $L \sim 10$kms, $v_o \sim 2\times10^4$cm/s, etc., we find $\frac{M}{m}\frac{\nu_{in}}{\nu_{en}}\frac{k_z^2}{k_\perp^2} \gg \left(\frac{\nu_{in}}{\Omega_i}\frac{1}{k_\perp L}\right)^2$, so that $\omega_R \sim 3.5 \times 10^{-2} s^{-1}$ and $\gamma \sim 10^{-3} sec^{-1}$, where $k_z/k \sim \left(M/m \frac{\nu_{in}}{\nu_{en}}\right)^{-\frac{1}{2}} \sim 10^{-2}$ was used with $\lambda_\perp \sim 300$ m. We note here that at E-region altitudes, an ambient electric field transverse to the magnetic field causes growth of perturbations due to the so-called $\underline{E} \times \underline{B}$ instability. The growth rate expression for this cross-field instability may be written as (Sudan et al., 1973)

$$\gamma \simeq \frac{\psi}{(1 + \psi)} \left(\frac{k_x \varepsilon}{k_\perp^2}\frac{\Omega_e}{\nu_{en}} \omega_R\right)$$

where $\omega_R \simeq \underline{k}\cdot\underline{v}_o/(1+\psi)$, $\psi \simeq \frac{\nu_{en}\nu_{in}}{\Omega_e\Omega_i}$, $\varepsilon \sim \frac{1}{n_o}\frac{dn_o}{dy}$, $v_o = \frac{cE_o}{B_o}$

(with E_o being the primary transverse electric field). Here temperature effects have been ignored for simplicity. The growth rate due to the $\underline{E} \times \underline{B}$ instability can be much higher depending upon the parameters. Thus, for $v_o \sim 100$ m/s, $L \sim 10$ km, $\gamma_{\underline{ExB}} \sim 10^{-1}s^{-1}$, in contrast to $\gamma_{cc} \sim 10^{-3}s^{-1}$ obtained above for the current convective instability.

C. Electromagnetic Effects

In this section we present results on the electromagnetic effects on the growth rate of the current convective instability. We consider the case in which temperature effects are ignored in addition to the ion and electron inertia. Instead of the electrostatic approximation, we use the full set of Maxwell's equations,

$$\nabla x\underline{E} = -\frac{1}{c}\frac{\partial \underline{B}}{\partial t} \; ; \; \nabla x\underline{B} = \frac{4\pi}{c}\underline{J} + \frac{1}{c}\frac{\partial \underline{E}}{\partial t}$$

The two equations above can be combined into one as,

$$\nabla^2\underline{E} - \nabla(\nabla\cdot\underline{E}) = \frac{4\pi}{c^2}\frac{\partial \underline{J}}{\partial t} \tag{23}$$

where the displacement current has been ignored for the low frequencies concerned. The perturbed current is

$$\underline{J} \simeq n_o e\,\underline{v}_i - n_o e\underline{v}_e - \tilde{n}ev_o\hat{z} \tag{24}$$

where the perturbed ion and electron fluid velocities are given by

$$\underline{v}_{\perp i} \simeq \frac{\nu_{in}}{\Omega_i^2}\frac{e\underline{E}_\perp}{M} + \frac{c\underline{E}_\perp x\hat{z}}{B_o} \; , \; v_{zi} \simeq \frac{eE_z}{M\nu_{in}} \tag{25}$$

and

$$\underline{v}_{\perp e} \simeq \frac{c}{B_o} \frac{(\omega - k_z v_o)}{\omega} \underline{E}_{\perp} x\hat{z} + \frac{c}{B_o} \frac{v_o}{\omega} \underline{k} x\hat{z}\, E_z, \quad v_{ze} \simeq - \frac{eE_z}{m\nu_e} \tag{26}$$

The rest of the perturbation analysis follows the procedure outlined in the preceding section. Use of (25) and (26) in (24) gives $\underline{J}$ which then is substituted back in eq. (23) to yield three component equations. The dispersion relation is obtained from the determinant of these three component equations and is, after some algebra,

$$\omega \left(\frac{\nu_e \nu_{in}}{\Omega_e \Omega_i} + \frac{k_z^2}{k_\perp^2} \right) - v_o k_z \frac{\nu_{in} \nu_{en}}{\Omega_e \Omega_i} - i \frac{v_o k_z}{\Omega_e} \frac{\nu_{en} \varepsilon}{k_\perp}$$

$$+ \frac{\omega_e^2}{c^2 k_\perp^2} \omega \frac{\nu_{in}}{\Omega_i} \left[- i \frac{\omega}{\Omega_e} + i \frac{v_o k_z}{\Omega_e} \frac{\nu_e \nu_{in}}{\Omega_e \Omega_i} - \frac{v_o k_z \varepsilon}{\Omega_e k_\perp} \frac{\nu_e}{\Omega_e} \right] \simeq 0. \tag{27}$$

The electrostatic limit can readily be obtained from eq. (27) by using the approximation $\frac{\omega_e}{ck_\perp} \ll 1$. In the opposite limit, $\frac{\omega_e}{ck_\perp} \gtrsim 1$, electromagnetic effects can be stabilizing or destabilizing depending on the parameters. However, for the auroral F-region situation they turn out to be negligibly small. The growth rate in this regime is, from eq. (27),

$$\gamma \approx \frac{v_o k_z \frac{\varepsilon}{k_\perp} \frac{\nu_e}{\Omega_e}}{\left(\frac{k_z^2}{k_\perp^2} + \frac{\nu_e \nu_i}{\Omega_e \Omega_i} \right)} - \frac{\omega_e^2}{c^2 k^2} \frac{\nu_{in} (v_o k_z)^2}{\Omega_e \Omega_i} \frac{\left(\frac{\varepsilon}{k_\perp} \right)^2 \left(\frac{\nu_e}{\Omega_e} \right)^2}{\left(\frac{k_z^2}{k_\perp^2} + \frac{\nu_e \nu_i}{\Omega_e \Omega_i} \right)^3} \tag{28}$$

For typical parameters, $\omega_e \sim 1.7\text{x}10^7$ rad s^{-1}, $\lambda_\perp \sim 1$ km, $\omega_e/ck_\perp \sim 10$, $\nu_{in}/\Omega_i \sim 2.5\text{x}10^{-4}$, $\nu_e/\Omega_e \sim 10^{-4}$, $\Omega_e \sim 5\text{x}10^6$, $v_o \sim 5\text{x}10^4$cm/s, $\lambda_z \sim 10^4$km, $L \sim 50$ km, the second term in eq. (28), though stabilizing, is $\sim 10^{-6}s^{-1}$ compared with the typical growth rates of the current convective instability in the auroral ionosphere which go as $\sim 10^{-3}s^{-1}$. Thus we find that in the

ionosphere the growth rate is unaffected by electromagnetic effects for the current convective instability. In a parameter regime where $\nu_{in}/\Omega_i > \varepsilon/k_\perp$, the electromagnetic effects on the current convective instability can be destabilizing, with the growth rate becoming,

$$\gamma \simeq \frac{\dfrac{v_o k_z}{\Omega_e}\dfrac{\epsilon}{k_\perp}\nu_e}{\left(\dfrac{k_z^2}{k_\perp^2}+\dfrac{\nu_e\nu_i}{\Omega_e\Omega_i}\right)} + \frac{\omega_e^2}{c^2k_\perp^2}\frac{\nu_{in}}{\Omega_i\Omega_e}\frac{(v_o k_z)^2\left(\dfrac{\nu_e\nu_e}{\Omega_e\Omega_i}\right)^2}{\left(\dfrac{k_z^2}{k_\perp^2}+\dfrac{\nu_e\nu_i}{\Omega_e\Omega_i}\right)^3} \qquad (28a)$$

This parameter regime may be satisfied at lower altitudes in the ionosphere.

III. DISCUSSION

We have presented in the previous sections the growth rate expressions in the case when ion-inertia is important (eqs. (12), (15), (16)); ions are highly collisional (eq. 22) *and in the case when the electromagnetic effects* are included (eq. (28)). This extends our previous analysis wherein we applied the current convective instability to the diffuse auroral situation ignoring ion-inertia, and using the electrostatic approximation in the analysis (eq. (13), (14)) and $\nu_{in} \ll \Omega_i$ [Ossakow and Chaturvedi, 1979].

For typical ionospheric F-region parameters at auroral latitudes at 350-400 km altitudes, we have $v_o \sim 500$ m/s, $m/M \sim 3\times10^{-5}$, $\nu_{ei} \sim 5\times10^2$, $\nu_{in} \sim 5\times10^{-2}$, $\varepsilon^{-1} \sim L \sim 50$ km, we see that from eq. (14), $\gamma_{CM} \sim 3\times10^{-3}\text{sec}^{-1}$ and similarly from eq. (16), $\gamma_{IM} \sim 3\times10^{-3}\text{sec}^{-1}$. We find that the growth rate remains the same order of magnitude in the collisional and inertial domains. At higher altitudes, where $\nu_{in} \lesssim 10^{-3}\text{s}^{-1}$, $\nu_{ei} \sim 30$, one finds that the growth rate for the same values of v_o and ε^{-1}, is $\sim 3\times10^{-3}\text{s}^{-1}$. Thus one can conclude that the current convective instability is applicable equally well to higher altitudes in the ionosphere where ν_{in} becomes smaller in contrast to

the collisional case presented in our previous work [Ossakow and Chaturvedi, 1979]. Thermal diffusion will damp out the modes at a rate proportional to $\sim c_s^2 (k_\perp^2 \frac{\nu_e}{\Omega_e \Omega_i} + \frac{k_z^2}{\nu_{in}})$ [Chaturvedi and Ossakow, 1979] which determines the threshold for the instability, and for a given set of parameters corresponding to instability, determines the cut-off scale lengths, below which the system is stable. For the present set of parameters, the stable (damped modes) wavelengths correspond to $\lambda_{\perp D} \leq 70$ m and $\lambda_{\parallel D} \leq 400$ km. In the limit of highly collisional ions, corresponding to E-region altitudes, the growth rate can be on the order of $\sim 10^{-3} \sec^{-1}$. However, the presence of transverse electric fields of moderate strengths (a few mv/m) would cause growth at a much faster rate due to the well-known $\underline{E} \times \underline{B}$ instability. The electromagnetic effects on the growth rate can be similarly estimated, and for the above set of parameters, for wavelengths ~ 1 km, they turn out to be a factor $\sim 10^{-3}$ smaller compared to the electrostatic growth rate, and are rather small to have any effect.

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